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Supersymmetry Breaking and Soft Terms in M-Theory

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Abstract

We investigate gaugino condensation in the framework of the strongly coupled heterotic $E_8 \times E_8$ string (M–theory). Supersymmetry is broken in a hidden sector and gravitational interactions induce soft breaking parameters in the observable sector. The resulting soft masses are of order of the gravitino mass. The situation is similar to that in the weakly coupled $E_8 \times E_8$ theory with one important difference: we avoid the problem of small gaugino masses which are now comparable to the gravitino mass.

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Phenomenological applications of string theory have mainly concentrated on the weakly coupled heterotic $E_8 \times E_8$ string. Models with realistic gauge groups and particle content have been constructed. Qualitatively the unification of gauge and gravitational couplings can be understood [1] with some moderate uncertainties in the relation between the string scale $M_{\rm string}$ and the grand unified scale $M_{\rm GUT}$. Supersymmetry is broken dynamically via gaugino condensation [2, 3] leading to a hidden sector supergravity model at low energies with specific predictions for the soft breaking parameters including problematically small gaugino masses [4]. At the moment we do not possess a convincing mechanism to resolve this problem.

Recently there have been attempts to study string theories in the region of intermediate and strong coupling. The strongly coupled version of the $E_8 \times E_8$ theory is believed to be an orbifold of 11–dimensional M–theory, an interval in $d = 11$ with $E_8 \times E_8$ gauge fields restricted to the two $d = 10$ dimensional boundaries respectively [5]. Applied to the question of unification [6] the following picture emerges: the GUT–scale $M_{\rm GUT} = 3 \times 10^{16}$ GeV is identified with $1/R$ where $V = R^6$ is the volume of compactified six–dimensional space. $\alpha_{\rm GUT} = 1/25$ and the correct value of the $d = 4$ reduced Planck mass $M_P = 2.4 \times 10^{18}$ GeV can be obtained by choosing the length of the $d = 11$ interval to be $R_{11} \approx 6R$. The fundamental mass scale of the $d = 11$ theory $M_{11} = \kappa^{-2/9}$ (with κ the $d = 11$ Einstein gravitational coupling) has to be chosen a factor 2 larger than $M_{\rm GUT}$ and at that scale $\alpha_{\rm string} = g^2_{\rm string}/4\pi$ is of order unity. This then represents a rather natural framework for the unification of coupling constants. The large compactification radius can be used to solve the strong CP problem [7, 8].

As in the weakly coupled theory, supersymmetry might be broken dynamically by gaugino condensation in the hidden E_8 on one boundary of space–time [9]. Gravitational interactions will play the role of messengers to the observable sector at the opposite boundary.

In this paper we shall discuss this mechanism in detail and compute the predictions for the low energy effective theory. We find results very similar to the situation in the $d = 10$ weakly coupled case, with one notable exception: *gaugino masses are of comparable size to the gravitino mass, thus solving the problem of small gaugino masses that occurred in the weakly coupled case.*

Let us first review gaugino condensation in the $d = 10$ weakly coupled $E_8 \times E_8$ theory. We shall start from the $d = 10$ effective field theory and go to $d = 4$ dimensions via the method of reduction and truncation explained in ref. [10]. In string theory compactified on an orbifold this would describe the dynamics of the untwisted sector. We retain the usual moduli fields S and T as well as matter fields C_i that transform nontrivially under the observable sector gauge group. In this approximation, the Kähler potential is given by [10, 11]

$$G = -\log(S + \bar{S}) - 3\log(T + \bar{T} - 2C_i\bar{C}_i) + \log|W|^2 \tag{1}$$

with superpotential

$$W(C) = d_{ijk}C_iC_jC_k \tag{2}$$

and the gauge kinetic function is given by the dilaton field

$$f = S\,. \tag{3}$$

We assume in the hidden sector the appearance of a gaugino condensate

$$\langle\chi\chi\rangle = \Lambda^3 \tag{4}$$

where Λ is the renormalization group invariant scale of the confining hidden sector gauge group. The gaugino condensate appears in the expression for the auxiliary components of the chiral superfields [12]

$$F_j = \left(G^{-1}\right)_j^k \left(\exp(G/2)G_k + \frac{1}{4} f_k(\chi\chi)\right) + \ldots \tag{5}$$

which are order parameters for supersymmetry breakdown. Minimizing the scalar potential we find $F_S = 0$, $F_T \neq 0$ and a vanishing cosmological constant. Supersymmetry is broken and the gravitino mass is given by

$$m_{3/2} = \frac{\langle F_T \rangle}{T + \bar{T}} \approx \frac{\Lambda^3}{M_P^2} \tag{6}$$

and $\Lambda = 10^{13}$ GeV would lead to a gravitino mass in the TeV – range. A first inspection of the soft breaking terms in the observable sector gives a disturbing result. They vanish in this approximation. Scalar masses are zero because of the no–scale structure in (1) (coming from the fact that we have only included fields of modular weight -1 under T–duality in this case) [4]. In a more general situation we would get scalar masses m_0 comparable to the gravitino mass $m_{3/2}$ and the above result $m_0 = 0$ is just an artifact of the chosen approximation at the classical level. Gaugino masses $m_{1/2}$ are given by

$$m_{1/2} = \frac{\frac{\partial f}{\partial S} F_S + \frac{\partial f}{\partial T} F_T}{2\mathrm{Re} f} \tag{7}$$

and with $f = S$ and $F_S = 0$ we obtain $m_{1/2} = 0$. One loop corrections will change this picture as can be seen already by an inspection of the Green–Schwarz anomaly cancellation counter terms, as they modify f at one loop. In the simple example of the so–called standard embedding with gauge group $E_6 \times E_8$ we obtain [11, 13]

$$f_6 = S + \epsilon T\,; \qquad\qquad f_8 = S - \epsilon T\,. \tag{8}$$

This dependence of f on T will via (7) lead to nonvanishing gaugino masses which, however, will be small compared to $m_{3/2}$ and m_0 since ϵT is considered a small correction to the classical result. This might be problematic when applied to the supersymmetric extension of the standard model.

Let us now move to the strongly coupled $E_8 \times E_8$ – M–theory. The effective action is given by [5]

$$\begin{aligned}
L \;=\; & \frac{1}{\kappa^2}\int d^{11}x\sqrt{g}\left[-\frac{1}{2}R - \frac{1}{2}\bar{\psi}_I \Gamma^{IJK} D_J\left(\frac{\Omega+\hat{\Omega}}{2}\right)\psi_K - \frac{1}{48} G_{IJKL}G^{IJKL}\right. \\
& \qquad -\frac{\sqrt{2}}{384}\left(\bar{\psi}_I \Gamma^{IJKLMN}\psi_N + 12\bar{\psi}^J\Gamma^{KL}\psi^M\right)\left(G_{JKLM} + \hat{G}_{JKLM}\right) \\
& \qquad \left. - \frac{\sqrt{2}}{3456}\epsilon^{I_1 I_2 \ldots I_{11}} C_{I_1 I_2 I_3} G_{I_4 \ldots I_7} G_{I_8 \ldots I_{11}}\right] \\
& + \frac{1}{2\pi(4\pi\kappa^2)^{2/3}}\int d^{10}x\sqrt{g}\left[-\frac{1}{4}F^a_{AB}F^{aAB} - \frac{1}{2}\bar{\chi}^a\Gamma^A D_A(\hat{\Omega})\chi^a\right. \\
& \qquad \left. -\frac{1}{8}\bar{\psi}_A \Gamma^{BC}\Gamma^A\left(F^a_{BC} + \hat{F}^a_{BC}\right)\chi^a + \frac{\sqrt{2}}{48}\left(\bar{\chi}^a\Gamma^{ABC}\chi^a\right)\hat{G}_{ABC11}\right].
\end{aligned} \tag{9}$$

Compactifying to $d = 4$ we obtain [6]

$$G_N = 8\pi\kappa_4^2 = \frac{\kappa^2}{16\pi^2 V\rho}\,, \qquad\qquad \alpha_{GUT} = \frac{(4\pi\kappa^2)^{2/3}}{2V} \tag{10}$$

with $V = R^6$ and $\pi\rho = R_{11}$. Fitting G_N and $\alpha_{GUT} = 1/25$ then gives $R_{11}M_{11} \approx 12$ and $M_{11}R \approx 2$ ($M_{11} \approx 6 \times 10^{16}$ GeV).

The rather large value of the $d = 4$ reduced Planck Mass $M_P = \kappa_4^{-1}$ is obtained as a result of the fact that R_{11} is large compared to R.

We now perform a compactification using the method of reduction and truncation as above. For the metric we write

$$g_{MN} = \begin{pmatrix} e^{-\gamma}e^{-2\sigma}g_{\mu\nu} & & \\ & e^{\sigma}\delta_{mn} & \\ & & e^{2\gamma}e^{-2\sigma} \end{pmatrix} \tag{11}$$

with $M, N = 1\ldots11$; $\mu,\nu = 1\ldots4$; $m,n = 5\ldots10$; $2R_{11} = 2\pi\rho = M_{11}^{-1}e^{\gamma}e^{-\sigma}$ and $V = e^{3\sigma}M_{11}^6$. At the classical level this leads to a Kähler potential as in (1)

$$K = -\log(\mathcal{S} + \bar{\mathcal{S}}) - 3\log(\mathcal{T} + \bar{\mathcal{T}} - 2C_i\bar{C}_i) \tag{12}$$

with

$$\mathcal{S} = \frac{2}{(4\pi)^{2/3}}\left(e^{3\sigma} \pm i24\sqrt{2}D\right)\,, \tag{13}$$

$$\mathcal{T} = \frac{\pi^2}{(4\pi)^{4/3}}\left(e^{\gamma} \pm i6\sqrt{2}C_{11}\right) \tag{14}$$

where D and C_{11} fields are defined by

$$\frac{1}{4!}e^{6\sigma}G_{11\lambda\mu\nu} = \epsilon_{\lambda\mu\nu\rho}\left(\partial^{\rho}D\right)\,, \tag{15}$$

$$C_{11i\bar{j}} = C_{11}\delta_{i\bar{j}} \tag{16}$$

and x^i ($x^{\bar{j}}$) is the holomorphic (antiholomorphic) coordinate of the Calabi–Yau manifold. The imaginary part of $\mathcal{S}$ (Im$\mathcal{S}$) corresponds to the model independent axion, and the gauge kinetic function is $f = \mathcal{S}$. This is very similar to the weakly coupled case. Before drawing any conclusion from these formulae, however, we have to discuss a possible obstruction at the one loop level. It can be understood from the mechanism of anomaly cancellation [6]. For the 3–index tensor field H in $d = 10$ supergravity to be well defined one has to satisfy $dH = \mathrm{tr}F_1^2 + \mathrm{tr}F_2^2 - \mathrm{tr}R^2 = 0$ cohomologically. In the simplest case of the standard embedding one assumes $\mathrm{tr}F_1^2 = \mathrm{tr}R^2$ locally and the gauge group is broken to $E_6 \times E_8$. Since in the M–theory case the two different gauge groups live on the two different boundaries of space–time such a cancellation point by point is no longer possible. We expect nontrivial vacuum expectation values (vevs) of

$$(dG) \propto \sum_i \delta(x^{11} - x_i^{11})\left(\mathrm{tr}F_i^2 - \frac{1}{2}\mathrm{tr}R^2\right) \tag{17}$$

at least on one boundary (x_i^{11} is the position of i–th boundary). In the case of the standard embedding we would have $\mathrm{tr}F_1^2 - \frac{1}{2}\mathrm{tr}R^2 = \frac{1}{2}\mathrm{tr}R^2$ on one and $\mathrm{tr}F_2^2 - \frac{1}{2}\mathrm{tr}R^2 = -\frac{1}{2}\mathrm{tr}R^2$ on

the other boundary. This might pose a severe problem since a nontrivial vev of G might be in conflict with supersymmetry ($G_{11ABC} = H_{ABC}$). The supersymmetry transformation law in $d = 11$ reads

$$\delta\psi_M = D_M\eta + \frac{\sqrt{2}}{288} G_{IJKL}\left(\Gamma_M^{IJKL} - 8\delta_M^I\Gamma^{JKL}\right)\eta + \dots \tag{18}$$

Supersymmetry will be broken unless e.g. the derivative term $D_M\eta$ compensates the nontrivial vev of G. Witten has shown [6] that such a cancellation can occur and constructed the solution in the linearized approximation (linear in the expansion parameter $\kappa^{2/3}$) which corresponds to the large T–limit in the weakly coupled theory[1]. The supersymmetric solution leads to a nontrivial dependence of the σ and γ fields with respect to x^{11}:

$$\frac{\partial\gamma}{\partial x^{11}} = -\frac{\partial\sigma}{\partial x^{11}} = \frac{\sqrt{2}}{24}\frac{\int d^6x\sqrt{g}\omega^{AB}\omega^{CD}G_{ABCD}}{\int d^6x\sqrt{g}} \tag{19}$$

where the integrals are over the Calabi–Yau manifold and ω is the corresponding Kähler form. A definition of our $\mathcal{S}$ and $\mathcal{T}$ fields in the four–dimensional theory would then require an average over the 11–dimensional interval. We would therefore write

$$\mathcal{S} = \frac{2}{(4\pi)^{2/3}}\left(e^{3\bar{\sigma}} \pm i24\sqrt{2}\bar{D}\right), \tag{20}$$

$$\mathcal{T} = \frac{\pi^2}{(4\pi)^{4/3}}\left(e^{\bar{\gamma}} \pm i6\sqrt{2}\bar{C}_{11}\right) \tag{21}$$

where bars denote averaging over the 11th dimension. It might be of some interest to note that the combination $\mathcal{S}\mathcal{T}^3$ is independent of x^{11} even before this averaging procedure took place.

$\exp(3\sigma)$ represents the volume of the six–dimensional compact space in units of M_{11}^{-6}. The x^{11} dependence of σ then leads to the geometrical picture that the volume of this space varies with x^{11} and differs at the two boundaries. In the given approximation, this variation is linear, and for growing R_{11} the volume on the E_8 side becomes smaller and smaller. At a critical value of R_{11} the volume will thus vanish and this will provide us with an upper limit on R_{11}. For the phenomenological applications we then have to check whether our preferred choice of R_{11} that fits the correct value of the $d = 4$ Planck mass[2] satisfies this bound. Although the coefficients are model dependent we find in general that the bound can be satisfied, but that R_{11} is quite close to its critical value. A choice of R_{11} much larger than $(5 \times 10^{15}\text{GeV})^{-1}$ is therefore not permitted.

This variation of the volume is the analogue of the one loop correction of the gauge kinetic function (8) in the weakly coupled case and has the same origin, namely a Green–Schwarz anomaly cancellation counterterm. In fact, also in the strongly coupled case we find corrections for the gauge coupling constants at the E_6 and E_8 side.

Gauge couplings will no longer be given by the (averaged) $\mathcal{S}$–field, but by that combination of the (averaged) $\mathcal{S}$ and $\mathcal{T}$ fields which corresponds to the $\mathcal{S}$–field before averaging

[1]For a discussion beyond this approximation in the weakly coupled case see ref. [14].

[2]With V depending on x^{11} we have to specify which values should be used in eq. (10). The appropriate choice in the expression for G_N is the average value of V while in the expression for α_{GUT} we have to use V evaluated at the E_6 border.

at the given boundary:

$$f_{6,8} = \mathcal{S} \pm \alpha \mathcal{T} \tag{22}$$

at the E_6 (E_8) side respectively[3]. The critical value of R_{11} will correspond to infinitely strong coupling at the E_8 side $\mathcal{S} - \alpha\mathcal{T} = 0$ (Notice the similarity to (8) in the weakly coupled limit). Since we are here close to criticality a correct phenomenological fit of $\alpha_{\rm GUT} = 1/25$ should include this correction $\alpha_{\rm GUT}^{-1} = \mathcal{S} + \alpha\mathcal{T}$ where $\mathcal{S}$ and $\alpha\mathcal{T}$ give comparable contributions. This is a difference to the weakly coupled case, where in $f = S + \epsilon T$ the latter contribution was small compared to S. Observe that this picture of a loop correction $\alpha\mathcal{T}$ to be comparable to the tree level result still makes sense in the perturbative expansion, since f does not receive further perturbative corrections beyond one loop [15, 16].

In a next step we are now ready to discuss the dynamical breakdown of supersymmetry via gaugino condensation in the strongly coupled M–theory picture. In analogy to the previous discussion we start investigating supersymmetry transformation laws in the higher–dimensional (now $d = 11$) field theory [9]:

$$\begin{aligned}
\delta\psi_A &= D_A\eta + \frac{\sqrt{2}}{288} G_{IJKL}\left(\Gamma_A^{IJKL} - 8\delta_A^I\Gamma^{JKL}\right)\eta \\
&\quad - \frac{1}{576\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\left(\bar\chi^a\Gamma_{BCD}\chi^a\right)\left(\Gamma_A^{BCD} - 6\delta_A^B\Gamma^{CD}\right)\eta + \ldots \qquad (23)\\
\delta\psi_{11} &= D_{11}\eta + \frac{\sqrt{2}}{288} G_{IJKL}\left(\Gamma_{11}^{IJKL} - 8\delta_{11}^I\Gamma^{JKL}\right)\eta \\
&\quad + \frac{1}{576\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\left(\bar\chi^a\Gamma_{ABC}\chi^a\right)\Gamma^{ABC}\eta + \ldots \qquad (24)
\end{aligned}$$

where gaugino bilinears appear in the right hand side of both expressions. It can therefore be expected that gaugino condensation breaks supersymmetry. Still the details have to be worked out. In the $d = 10$ example, the gaugino condensate and the three–index tensor field H contributed to the scalar potential in a full square. This lead to a vanishing cosmological constant as well as the fact that $F_S = 0$ at the classical level. Hořava has observed [9] that a similar mechanism might be in operation in the $d = 11$ theory

$$\begin{aligned}
-\frac{1}{12\kappa^2}\int_{M^{11}} d^{11}x\sqrt{g}G_{ABC11}^2 &+ \frac{\sqrt{2}}{24(4\pi)^{5/3}\kappa^{4/3}}\int_{M^{10}} d^{10}x\sqrt{g}G_{11ABC}\left(\bar\chi^a\Gamma^{ABC}\chi^a\right) \\
&- \frac{\delta(0)}{96(4\pi)^{10/3}\kappa^{2/3}}\int_{M^{10}} d^{10}x\sqrt{g}\left(\bar\chi^a\Gamma^{ABC}\chi^a\right)^2 \qquad (25)\\
= -\frac{1}{12\kappa^2}\int_{M^{11}} d^{11}x\sqrt{g}&\left(G_{ABC11} - \frac{\sqrt{2}}{16\pi}\left(\frac{\kappa}{4\pi}\right)^{2/3}\delta(x^{11})\bar\chi^a\Gamma_{ABC}\chi^a\right)^2 .
\end{aligned}$$

After a careful calculation this leads to a vanishing variation $\delta\psi_A = 0$. In our model (based on reduction and truncation) we can now compute these quantities explicitly. We assume gaugino condensation to occur at the E_8 boundary

$$\langle\bar\chi^a\Gamma_{ijk}\chi^a\rangle = \Lambda^3\epsilon_{ijk} \tag{26}$$

[3] With the normalization of the $\mathcal{T}$ field as in (21), α is a quantity of order 1.

where $\Lambda < M_{\rm GUT}$ and ϵ_{ijk} is the covariantly constant holomorphic 3–form. This leads to a nontrivial vev of G_{11ABC} at this boundary and supersymmetry is broken[4]. At that boundary we obtain $F_{\mathcal{S}} = 0$ and $F_{\mathcal{T}} \neq 0$ as expected from the fact that the component ψ_{11} of the 11–dimensional gravitino plays the role of the goldstino.

In the effective $d = 4$ theory we now have to average over the 11th dimension leading to

$$\langle F_{\mathcal{T}} \rangle \approx \frac{1}{2}\mathcal{T}\frac{\int dx^{11}\delta\psi_{11}}{\int dx^{11}} \tag{27}$$

as the source of SUSY breakdown. This will then allow us to compute the size of supersymmetry breakdown on the observable E_6 side. Gravitational interactions play the role of messengers that communicate between the two boundaries. This effect can be seen from (10): large R_{11} corresponds to large M_P and $\langle F_{\mathcal{T}} \rangle$ gives the effective size of SUSY breaking on the E_6 side ($R_{11} \to \infty$ implies $M_P \to \infty$). The gravitino mass is given by

$$m_{3/2} = \frac{\langle F_{\mathcal{T}} \rangle}{\mathcal{T} + \bar{\mathcal{T}}} \approx \frac{\Lambda^3}{M_P^2} \tag{28}$$

(similar to (6) in the weakly coupled case) and we expect this to represent the scale of soft supersymmetry breaking parameters in the observable sector. These soft masses are determined by the coupling of the corresponding fields to the goldstino multiplet. As we have seen before, we cannot compute the scalar masses reliably in our approximation: $m_0 = 0$ because of the no–scale structure that appears as an artifact of our approximation. Fields of different modular weight will receive contribution to m_0 of order $m_{3/2}$. For the mass of a field C we have [17, 4]

$$m_0^2 = m_{3/2}^2 - F^i \bar{F}^{\bar{j}} \frac{Z_{i\bar{j}} - Z_i Z^{-1} \bar{Z}_{\bar{j}}}{Z} \tag{29}$$

where $i, j = \mathcal{S}, \mathcal{T}$ and Z is the moduli dependent coefficient of $C\bar{C}$ term appearing in the Kähler potential. Scalars of modular weight -1 will become massive through radiative corrections. This then leads to the expectation that $m_{3/2}$ should be in the TeV–region and $\Lambda \approx 10^{13}$ GeV [5]. So far this is all similar to the weakly coupled case.

An important difference appears, however, when we turn to the discussion of observable sector gaugino masses (7). In the weakly coupled case they were zero at tree level and appeared only because of the radiative corrections at one loop (8). As a result of this small correction, gaugino masses were expected to be much smaller than $m_{3/2}$. In the strongly coupled case the analog of (7) is still valid

$$m_{1/2} = \frac{\frac{\partial f_6}{\partial \mathcal{S}} F_{\mathcal{S}} + \frac{\partial f_6}{\partial \mathcal{T}} F_{\mathcal{T}}}{2{\rm Re} f_6} \tag{30}$$

and the 1–loop effect is encoded in the variation of the σ and γ fields from one boundary to the other. Here, however, the loop corrections are sizable compared to the classical

[4] One might speculate that a nontrivial vev of $D_A\eta$ might be operative here as in the case without gaugino condensation (see discussion after eq. (18)). However, the special values of $H_{ijk} \propto \epsilon_{ijk}$ necessary to cancel the contribution of the gaugino condensate do not permit such a mechanism (see footnote 6 in ref. [6]).

[5] In realistic models E_8 is broken and Λ is adjusted by model building.

result because of the fact that R_{11} is close to its critical value. As a result we expect observable gaugino masses of the order of the gravitino mass. The problem of the small gaugino masses does therefore not occur in this situation. Independent of the question whether $F_{\mathcal{S}}$ or $F_{\mathcal{T}}$ are the dominant sources of supersymmetry breakdown, the gauginos will be heavy of the order of the gravitino mass. The exact relation between the soft breaking parameters m_0 and $m_{1/2}$ will be a question of model building. If in some models $m_0 \ll m_{1/2}$ this might give a solution to the flavor problem. The no–scale structure found above might be a reason for such a suppression of m_0. As we have discussed above this structure, however, is an artifact of our simplified approximation and does not survive in perturbation theory. At best it could be kept exact (but only for the fields with modular weight -1) in the $R_{11} \to \infty$ limit. The upper bound on R_{11} precludes such a situation. With observable gaugino masses of order $m_{3/2}$ we also see that $m_{3/2}$ cannot be arbitrarily large and should stay in the TeV – range.

In recent months several other groups have studied similar questions in detail [18, 19, 20]. The discussion was explicitly done at the classical level. Some conclusions different from ours (concerning large values of $m_{3/2}$ and/or R_{11}) can only be obtained in that approximation. The one loop corrections, however, require $R_{11} < R_{\rm critical}$ as well as $m_{3/2}$ in the TeV – range.

The picture of supersymmetry breakdown in the M–theoretic limit therefore seems very promising. It is very similar to the weakly coupled case, but avoids the problem of the small gaugino masses.

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